Cabibbo-favored nonleptonic decays of charmed baryons

Q. P. Xu and A. N. Kamal
Theoretical Physics Institute and Department of Physics, University of Alberta, Edmonton, Alberta, Canada T6G 2J1
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We address several important issues in nonleptonic decays of charmed baryons. These include the treatment of the baryon matrix element of the effective weak Hamiltonian and the factorization in the quark-decay part of the amplitude. We also introduce an on-shell correction to the soft-pion limit of the current-algebra result for S waves. This correction leads to significant changes in the decay rate and, in particular, the asymmetry parameter. A calculation of Cabibbo-favored $B_c \to B + P(0^-)$ decays of Λ_c^+, Ξ_c^{0A}, and Ξ_c^{+A} (where A indicates asymmetry in the noncharmed quarks) has been carried out. "Good candidates" for experimental tests have been suggested.

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I. INTRODUCTION

As more and more experimental data on charmed baryons become available [1–7], the theoretical study of nonleptonic decays of charmed baryons has become very important. Unlike the hyperon decays, there are more decay channels available to charmed baryons providing a rich testing ground for the standard model. This is quite similar to the situation in charmed-meson decays as compared to that in strange-meson decays. Theoretically, charmed-baryon decays are more difficult to handle than hyperon and charmed-meson decays. In hyperon decays, the W-exchange or equivalently the pole term induced by the strong quark-quark correlation is the dominant contribution while factorization plays a supplementary role [8]. In charmed-meson decays, the factorization contribution dominates over W exchange, i.e., weak annihilation [9,10]. Together with a proper treatment of final-state interactions (FSI) [9,11,12], the factorization model has been quite successful in explaining much of the charmed-meson-decay data. In charmed baryon decays naive expectation would lead us to believe that factorization would also make the dominant contribution due to the large amount of energy emitted but, as we will show later, this is not the case. Both factorization and pole contributions are equally important, though one could dominate over the other depending on the decay mode. The situation may, however, be further complicated by the FSI effect.

Theoretical studies of charmed-baryon decays began a long time ago [13–16] and currently are under intensive investigation [17–20]. We have referred here only to dynamical models. In most such studies the treatment is based on current-algebra techniques. Our basic assumption is that factorization and the pole term are the two possible contributions to the decay amplitude. However, we will emphasize several important issues in the treatment of charmed-baryon decays which can completely change the decay behavior in many channels. We first point out the very significant difference due to different treatment of the factorization contribution. Although quite well known in charmed-meson decays [9], this point has not yet been made very clear in charmed baryon decays. Firstly, we will use in our calculation the so-called "new factorization" which is essential in understanding charmed meson decays [9]. Secondly, for the baryon matrix element of the weak Hamiltonian we will not use the value obtained from the corresponding hyperon matrix element and SU(4) symmetry as in [17–19] but calculate this quantity using the quark-quark correlation mechanism [8,21]. Thirdly, the S-wave amplitude has always been represented by the commutator term in the soft limit; however, away from the soft limit there arises an additional contribution which we call the "on-shell correction." This on-shell correction is important in that it alters the S-wave amplitude and thus has a testable effect in the decay rate and, especially, in the decay asymmetry parameter α. In this paper we will concentrate on Cabibbo-favored $B_c \to B(\frac{1}{2}^+) + P(0^-)$ weak decays of Λ_c^+, Ξ_c^{0A}, and Ξ_c^{+A} charmed baryons (superscripts A and S indicate antisymmetry and symmetry in the noncharmed quarks). However, we will also suggest several other types of decays which are very interesting due to their simplicity.

This paper is organized as follows: We give the general formalism in Sec. II. A discussion on factorization is also included here. In Sec. III we treat the baryon matrix element of the weak Hamiltonian and in Sec. IV we show how to determine the on-shell correction. This is followed by a numerical calculation and a discussion in Sec. V. We end with conclusions in Sec. VI.

II. GENERAL FORMALISM

For Cabibbo-favored process ($\Delta S = \Delta C = 1$), the relevant effective weak Hamiltonian including QCD short-range effects has the form

$$H_w^{\text{eff}}(\Delta S = \Delta C = 1) = \frac{G_F V_{ud} V_{cs}}{\sqrt{2}}(c_1 Q_1 + c_2 Q_2) , \qquad (1)$$

where

$$Q_1 = [\bar{s}\gamma_\mu(1-\gamma_5)c][\bar{u}\gamma_\mu(1-\gamma_5)d] , \qquad (2)$$

$$Q_2 = [\bar{u}\gamma_\mu(1-\gamma_5)c][\bar{s}\gamma_\mu(1-\gamma_5)d] . \qquad (3)$$

The Wilson coefficients are chosen to be [22]

$$c_1=1.2,\quad c_2=-0.5\ . \tag{4}$$

The amplitude for the decay $B_c \to B+P(0^-)$ is defined by [23]

$$\langle B,P|H_w^{\rm eff}|B_c\rangle = i\bar{u}_B(A-\gamma_5 B)u_{B_c}\ , \tag{5}$$

where B and B_c are the ground-state charmed and noncharmed $\frac{1}{2}^+$ baryons. The various contributions to the decay amplitudes are discussed in the following.

A. The factorization contribution

The simplest part of the decay amplitude is obtained by treating the quark current in the weak Hamiltonian as an interpolating hadron field and therefore one of the currents in Q_1 or Q_2 of (1) directly generates a pseudoscalar meson. We call this the factorization contribution. The result of factorization depends on the scale at which it is performed. From the structure of the effective weak Hamiltonian, only channels $B_c \to B+P(0^-)$ with $P=\pi^+$ or $\bar{K}^0$ can receive a factorization contribution. At first sight one might expect factorization to be the most important contribution due to the large amount of energy emitted in charmed-baryon decays. For decays with π^+ emission, the corresponding amplitudes are given by

$$A^{\rm fact}=\frac{G_F V_{ud}V_{cs}}{\sqrt{2}}a_1(M_{B_c}-M_B)f_\pi f^V_{B_cB}\ , \tag{6}$$

$$B^{\rm fact}=-\frac{G_F V_{ud}V_{cs}}{\sqrt{2}}a_1(M_{B_c}+M_B)f_\pi g^A_{B_cB}\ . \tag{7}$$

For $\bar{K}^0$ emission a_1 and f_π should be replaced by a_2 and f_K. The coefficients a_1 and a_2 are given by

$$a_1=c_1+\xi c_2,\quad a_2=c_2+\xi c_1\ . \tag{8}$$

The terms ξc_1 and ξc_2 arise from Fierz reordering of the operators Q_1 and Q_2. One crucial question in the treatment of the factorization contribution is how to handle the Fierz-reordering term. If one keeps only the color-singlet current induced by Fierz reordering and neglects the color-octet current, one gets $\xi=1/N_c$. Earlier treatments of the factorization contribution had assumed $\xi=\frac{1}{3}$ [13–16]. We refer to this as "old factorization." The so-called "new factorization," which neglects the Fierz-reordering term, is equivalent to setting $\xi\to 0$. The effect of this so-called "new factorization" is to enhance the amplitudes involving $\bar{K}^0$ in the final state, leaving the amplitudes involving π^+ emission essentially unaltered. This is analogous to an enhancement of the $D^0\to\pi^0\bar{K}^0$ amplitude in D^0 decays leaving the $D^0\to\pi^+K^-$ amplitude essentially uneffected [9]. Additional support for this treatment comes from the $1/N_c$ expansion method [24] and the QCD sum-rule calculations [25]. We will therefore apply the new factorization here.

The axial-vector and vector current form factors of baryons in (6) and (7) are defined with the following first-order parametrization:

$$\langle B|V_\mu|B_c\rangle = f^V_{BB_c}\bar{u}_B\gamma_\mu u_{B_c};\ \langle B|A_\mu|B_c\rangle = g^A_{BB_c}\bar{u}_B\gamma_\mu\gamma_5 u_{B_c}\ . \tag{9}$$

For numerical values of form factors $f^V_{BB_c}$ and $g^A_{BB_c}$, we use the results of a recent calculation by Avila-Aoki *et al.* using the bag model and a dipole q^2 dependence [26]. They have calculated $(c\to s)$-type form factors. For example, the $\Lambda_c^+\to\Lambda$ form factors are given by

$$\begin{aligned} f^V_{\Lambda\Lambda_c^+}(q) &= \frac{f^V_{\Lambda\Lambda_c^+}(0)}{(1-q^2/m^2_{1^-})^2}\ ,\\ g^A_{\Lambda\Lambda_c^+}(q) &= \frac{g^A_{\Lambda\Lambda_c^+}(0)}{(1-q^2/m^2_{1^+})^2}\ , \end{aligned} \tag{10}$$

where $f^V_{\Lambda\Lambda_c^+}(0)=-0.46$ and $g^A_{\Lambda\Lambda_c^+}(0)=-0.50$ and the pole masses are $m_{1^-}=2.1$ GeV and $m_{1^+}=2.5$ GeV. Form factors induced by the $c\to u$ current have not been given in [26] but they can be related through SU(3) symmetry to the calculated $c\to s$ form factors; for instance, $[f^V_{P\Lambda_c^+}(0),\ g^A_{P\Lambda_c^+}(0)]=\sqrt{\frac{3}{2}}[f^V_{\Lambda\Lambda_c^+}(0),\ g^A_{\Lambda\Lambda_c^+}(0)]$. The corresponding current pole masses are $m_{1^-}=2.0$ GeV and $m_{1^+}=2.4$ GeV.

B. The pole contribution

Besides describing a direct emission of mesons, the weak Hamiltonian will also induce a mixing of hadron fields in any effective hadron Hamiltonian. For instance, the Λ_c^+ field will mix through $H_\omega^{\rm eff}$ in (1) with $\Sigma^+(\frac{1}{2}^+)$ and $\Sigma^+(\frac{1}{2}^-)$ carrying the Σ^+ flavor quantum numbers. The corresponding charmed baryon decays are then obtained from the strong-interaction baryon-meson vertex. This is just the picture described by the s-channel pole term. A similar picture is described by the u-channel pole term. The pole contribution plays a dominant role in hyperon decays, but in charmed-baryon decays studied here, a naive expectation is that it is not as important as factorization. Obviously, the pole contribution is very important in channels without π^+ or $\bar{K}^0$ emission where the factorization term vanishes.

Formulas for $\frac{1}{2}^+$ and $\frac{1}{2}^-$ pole terms are well known [8,27,28]. For P waves in $B_c\to B+P(0^-)$, the expression reads

$$B^{\rm pole}=\frac{(M_B+M_{B_c})}{f_P}\left[\frac{g^A_{BB'}h^+_{B'B_c}}{(M_{B_c}-M_{B'})}+\frac{h^+_{BB_c'}g^A_{B_c'B_c}}{(M_B-M_{B_c'})}\right], \tag{11}$$

where f_P can be either f_π or f_K. The s- and u-channel poles B' and B_c' are ground-state $\frac{1}{2}^+$ noncharmed and charmed baryons, respectively. The baryon matrix element of the parity-conserving part of $H_w^{\rm eff}$ has the general Lorentz structure

$$\langle B_1(\tfrac{1}{2}^+)|H_w^{\rm eff(pc)}|B_2(\tfrac{1}{2}^+)\rangle = h^+_{B_1B_2}\bar{u}_{B_1}u_{B_2}\ . \tag{12}$$

It should be noted that (11) is in fact a subtracted pole formula, i.e., it is obtained by subtracting from $B^{\rm pole}(q)$ its $q\to 0$ analytic continuation to ensure that $B^{\rm pole}\to 0$ in the soft limit $q\to 0$. (Recall that in the soft limit the P-wave amplitude vanishes.) The form factors in (11) are of two types, those between noncharmed baryons $g^A_{BB'}$ and those between charmed baryons $g^A_{B'_cB_c}$ (but $\Delta C=0$). For the former, we use the usual SU(3) parametrization with

$$D+F=1.25,\quad D/F\approx 1.8\ . \tag{13}$$

The D/F ratio is taken from a fit to hyperon semileptonic decays [29]. Form factors $g^A_{B'_cB_c}$ can also be given by D and F through SU(4) symmetry. For instance, $g^A_{\Sigma^0_c\Lambda^+_c}=g^A_{\Sigma^-\Lambda}=\sqrt{\tfrac{2}{3}}D$, if one makes the substitution $c\to s$. The SU(4) relation for $g^A_{B'_cB_c}$ is expected to be a good approximation since the transition involved here is $\Delta C=0$ and the baryon wave function mismatch in the overlap integral is small.

The S-wave pole formula is (see [8] for the method)

$$A^{\rm pole}=-\frac{1}{f_P}\left[\frac{(M^*_{B'}-M_B)}{(M^*_{B'}-M_{B_c})}E^A_{BB'}h^-_{B'B_c}-\frac{(M^*_{B'_c}-M_{B_c})}{(M^*_{B'_c}-M_B)}h^-_{BB'_c}E^A_{B'_cB_c}\right], \tag{14}$$

where $M^*_{B'}$ and $M^*_{B'_c}$ are the masses of excited $\frac{1}{2}^-$ noncharmed baryons $B'(\frac{1}{2}^-)$ and charmed baryons $B'_c(\frac{1}{2}^-)$, coupled to the ground-state $\frac{1}{2}^+$ baryons through the parity-violating $H_w^{\rm eff(pv)}$. The corresponding baryon matrix element of $H_w^{\rm eff}$ is defined by

$$\langle B_1(\tfrac{1}{2}^+)|H_w^{\rm eff(pv)}|B_2(\tfrac{1}{2}^-)\rangle=h^-_{B_1B_2}\bar u_{B_1}u_{B_2}\ . \tag{15}$$

The axial-vector current matrix element between $\frac{1}{2}^+$ and $\frac{1}{2}^-$ baryons is given by

$$\langle B_1(\tfrac{1}{2}^+)|A_\mu|B_2(\tfrac{1}{2}^-)\rangle=E_{B_1B_2}\bar u_{B_1}\gamma_\mu u_{B_2}\ . \tag{16}$$

The pole contribution $A^{\rm pole}$ should give, in the soft limit $q\to 0$, the well-known commutator term derived from current algebra [8], i.e.,

$$A^{\rm pole}(q\to 0)=A^{\rm soft}=-\frac{1}{f_P}(E^A_{BB'}h^-_{B'B_c}-h^-_{BB'_c}E^A_{B'_cB_c}) \tag{17}$$

and

$$A^{\rm soft}=A^{\rm com}=-\frac{1}{f_P}\langle B|[I^5,H_w^{\rm eff(pv)}]|B_c\rangle\ . \tag{18}$$

Thus,

$$A^{\rm pole}=A^{\rm com}+\frac{(M_{B_c}-M_B)}{f_P}\left[\frac{E^A_{BB'}h^-_{B'B_c}}{(M_{B_c}-M^*_{B'})}-\frac{h^-_{BB'_c}E^A_{B'_cB_c}}{(M^*_{B'_c}-M_B)}\right]. \tag{19}$$

Thus the $\frac{1}{2}^-$ pole term contains the soft-limit result, i.e., the commutator term, and a second term which we call the on-shell correction. This correction obviously vanishes in the soft limit, as can be seen from (19). In previous studies of charmed baryon decays, this on-shell correction has not been included. The importance of this correction can be estimated, roughly, as follows. If $M^*_{B'_c}(\frac{1}{2}^-)\approx M_{B_c}+0.3$ GeV and $M^*_{B'}(\frac{1}{2}^-)\approx M_B+0.6$ GeV, then $(M_{B_c}-M_B)/(M_{B_c}-M^*_{B'})\approx 2$ and $(M_{B_c}-M_B)/(M^*_{B'_c}-M_B)\approx 0.7$, indicating a possible large on-shell correction. In hyperon decays a 30% on-shell correction to the soft limit is obtained and this correction is a very important step in obtaining a successful description of hyperon nonleptonic decays [8].

Finally, the total amplitude is given, as usual, by

$$A=A^{\rm fact}+A^{\rm pole},\quad B=B^{\rm fact}+B^{\rm pole}\ . \tag{20}$$

III. BARYON MATRIX ELEMENT OF $H_w^{\rm eff}$

In this section we will discuss how to calculate the so-far unknown baryon matrix element of $H_w^{\rm eff}$ defined in (12).

In some recent studies of charmed baryon decays [17–19], SU(4) symmetry has been used to get baryon matrix elements of the weak Hamiltonian, say $h^+_{\Sigma^+\Lambda_c^+}$, from hyperon nonleptonic decays

$$h^+_{\Sigma^+\Lambda^+_c}=\frac{1}{\sqrt{6}}\cot\theta_C\,h^+_{P\Sigma^+}\ , \tag{21}$$

where θ_C is the Cabibbo angle. The above relation follows if a $c\to s$ substitution is made. However, it is quite problematic to use SU(4) symmetry here since the large mass difference between c and s quarks is expected to lead to a large mismatch in the baryon wave functions used in the overlap integral leading to the real baryon matrix element of $H_w^{\rm eff}(\Delta C=1)$ being smaller than that given by (21). This influences the importance of the pole term.

To calculate the baryon matrix element of $H_w^{\rm eff}$ in (12), we will apply the quark-quark correlation mechanism which was found to be crucial in understanding the nonleptonic weak decays of strange particles [21]. With this mechanism, the longstanding $\Delta I=\frac{1}{2}$ rule was well explained [21] and a coherent and a very successful description of K decays [30], the K_L-K_S mass difference [31], and hyperon nonleptonic decays [8] was accomplished. This mechanism has also been found to be important in understanding the inclusive charm-decay rates [21] and has been applied to the baryonic decays of B decays [32]. To be specific, we give a simple picture of this mechanism in the baryon matrix element of $H_w^{\rm eff}$ relevant to our case.

The reader is referred to the original papers for details. The effective weak Hamiltonian can be written in the following form after a Fierz transformation [21]:

$$H_w^{\text{eff}}(\Delta S=\Delta C=1)$$
$$=\frac{G_F V_{ud} V_{cs}}{\sqrt{2}}[c_-(su)^+_{3^*}(cd)_{3^*}+c_+(su)^+_6(cd)_6] \quad (22)$$

with $c_-=c_1-c_2, c_+=c_1+c_2$. $(su)_{3^*}=\epsilon_{ijk}\overline{s}^c_j(1-\gamma_5)u_k$ ($\overline{s}^c=s^T C$, C =charge conjugation) is a local pseudoscalar and scalar color-antitriplet diquark current (i,j,k are color indices). Equation (22) shows that the operator $Q_-(=Q_1-Q_2)$ is the product of color-antitriplet diquark currents while $Q_+(=Q_1+Q_2)$ the product of color-sextet diquark currents. In baryons, two quarks can only be in a color-antitriplet state. At short distances the force between two quarks is attractive in the color **3*** state and repulsive in the color **6** state. This is already reflected in the enhancement of c_- and suppression of c_+. The force between two quarks in the color **3*** state at large distances is also strongly attractive, as can be seen from the very existence of baryons. The concept of diquarks, a quasibound state, as a result of such quark-quark correlation, has wide implications in hadron physics and spectroscopy [33]. In the quark-quark correlation mechanism, the weak effective Hamiltonian in (22), when acting between two baryons, simply annihilates a scalar or pseudoscalar $(cd)_{3^*}$ diquark from the initial baryon and then creates a corresponding $(us)_{3^*}$ diquark in the final baryon. The measure of the annihilation and creation of diquarks through the diquark current in (22) is given by the "diquark decay constant"

$$\langle 0|\epsilon_{ijk}\overline{s}^c_j\gamma_5 u_k|(su)^{0+}_l\rangle=\sqrt{\tfrac{2}{3}}\delta_{il}g_{su}\ ,$$
$$\langle 0|\epsilon_{ijk}\overline{c}^c_j\gamma_5 d_k|(cd)^{0+}_l\rangle=\sqrt{\tfrac{2}{3}}\delta_{il}g_{cd}\ . \quad (23)$$

In the above 0^+ indicates a scalar diquark. Similarly, the decay constant of a pseudoscalar diquark can also be defined. The diquark decay constants in (23) have been given in [34,21]:

$$g_{su}=0.215\pm0.013\ \text{GeV}^2\ ,$$
$$g_{cd}=0.35\pm0.05\ \text{GeV}^2\ . \quad (24)$$

With these decay constants it is then possible to calculate the baryon matrix element of H_w^{eff}. Because of the quark-quark correlation mechanism, the result is quite insensitive to the details of the wave functions involved. Such a calculation has already been carried out for hyperon nonleptonic decays [8,35]. It was found that

$$h^+_{P\Sigma^+}=\sqrt{2}(d-f)\simeq 1.1\times10^{-7}\ \text{GeV}$$

or (25)

$$d=-f\simeq 0.38\times10^{-7}\ ,$$

where d and f are the usual SU(3) parameters defined by $h^+_{B_iB_j}=2\sqrt{2}(if_{i6j}f+d_{i6j}d)$, B_i and B_j being octet baryons. Here for the charmed-baryon matrix element we follow the same procedure [8,35] and use (24). We find

$$h^+_{\Sigma^+\Lambda_c^+}\simeq 0.8\times10^{-7}\ \text{GeV}\ . \quad (26)$$

Comparison of (26) with (21) and (25) shows that the SU(4) relation is badly broken. (Our calculated value (26) is quite consistent with the bag-model result $h^+_{\Sigma^+\Lambda_c^+}\approx 1.0\times10^{-7}$ GeV with $c_-=1.96$, calculated in the earlier studies of charmed baryon decays [14–16].) We use (26) in our calculations. Other charmed-baryon matrix elements can be obtained from (26) through SU(3) symmetry. Here we list the relations relevant to our later calculation:

$$h^+_{\Sigma^+\Lambda_c^+}=h^+_{\Sigma^+\Sigma_c^+}/\sqrt{3}=-h^+_{\Sigma^0\Sigma_c^0}/\sqrt{3}$$
$$=h^+_{\Lambda\Sigma_c^0}=-h^+_{\Xi^0\Xi_c^{0A}}=h^+_{\Xi^0\Xi_c^{0S}}/\sqrt{3}\ . \quad (27)$$

IV. THE $\frac{1}{2}^-$ POLE TERM AND ON-SHELL CORRECTION

We start with a discussion of the S-wave pole term. The soft limit, i.e., the commutator term is completely determined by Eq. (18) and the baryon matrix element of H_w^{eff} in (26) discussed in the preceding section. For the on-shell correction, the pole baryons are the excited $\frac{1}{2}^-$ counterparts of those $\frac{1}{2}^+$ baryon ground states which contribute to the P-wave pole term. In addition, the excited $\frac{1}{2}^-$ SU(4) 4-plet baryons, in our notation Λ_c^{+1}, $\Xi_c^{0\,1}$, Ξ_c^{+1}, Λ^1 [corresponding to the excited $\frac{1}{2}^-$ SU(3) singlet Λ^1], can also contribute to the S-wave pole terms. The on-shell correction in (19) involves the form factors and matrix elements with $\frac{1}{2}^+\to\frac{1}{2}^-$ transition. We denote the analogues of the D and F parameters, Eq. (13), for $\frac{1}{2}^+\to\frac{1}{2}^-$ baryon-baryon form factors, in the following, by D^- and F^-. The symmetry relations among $\frac{1}{2}^+\to\frac{1}{2}^-$ baryon form factors are the same as the corresponding ones among the $\frac{1}{2}^+\to\frac{1}{2}^+$ ground-state baryon form factors. Similarly the relation in (27) is also true for $\frac{1}{2}^+\to\frac{1}{2}^-$ baryon matrix elements. To calculate the on-shell correction, we go back to (17) and (18). The commutator term is the soft limit of the $\frac{1}{2}^-$ pole term. From the commutator term (18) and relation (17) we can obtain restrictions on the product of baryon form factor and matrix elements of the Hamiltonian:

$$(D^--F^-)h^-_{\Sigma^+\Lambda_c^+}=h^+_{\Sigma^+\Lambda_c^+}\ , \quad (28)$$

$$e=\tfrac{2}{3}D^-h^-_{\Sigma^+\Lambda_c^+}\ , \quad (29)$$

where e is proportional to the product $h^-_{BB_c'}E^A_{B_c'B_c}$ coming from the excited SU(4) 4-plet $\frac{1}{2}^-$ baryon poles, i.e., $B_c'=\Lambda_c^{+1}$, $\Xi_c^{0\,1}$, Ξ_c^{+1}, or Λ^1. The products above, not the individual form factor or baryon matrix element, determine the on-shell correction [36]. If we define

$$x=\frac{(F^-+D^-)}{(F^--D^-)} \quad (30)$$

then we can write the $\frac{1}{2}^-$ pole contributions in terms of the commutator term. Here we give the resulting formulas for Λ_c^+ decays:

$$A^{\text{pole}}(\Lambda_c^+\to\Lambda\pi^+)=\left[\frac{h^+_{\Sigma^+\Lambda_c^+}}{f_\pi}\right](M_{\Lambda_c^+}-M_\Lambda)\left[\frac{(1-x)}{\sqrt{6}(M_{\Lambda_c^+}-M^*_{\Sigma^+})}-\frac{(1-x)}{\sqrt{6}(M^*_{\Sigma_c^0}-M_\Lambda)}\right], \tag{31}$$

$$A^{\text{pole}}(\Lambda_c^+\to P\bar{K}^0)=\left[\frac{-h^+_{\Sigma^+\Lambda_c^+}}{f_K}\right]\left\{1-\frac{(M_{\Lambda_c^+}-M_P)}{(M_{\Lambda_c^+}-M^*_{\Sigma^+})}\right\}, \tag{32}$$

$$A^{\text{pole}}(\Lambda_c^+\to\Sigma^0\pi^+)=\left[\frac{-\sqrt{2}h^+_{\Sigma^+\Lambda_c^+}}{f_\pi}\right]\left[1-(M_{\Lambda_c^+}-M_{\Sigma^0})\left[\frac{(1+x)}{2(M_{\Lambda_c^+}-M^*_{\Sigma^+})}+\frac{(1-x)}{2(M^*_{\Sigma_c^0}-M_{\Sigma^0})}\right]\right], \tag{33}$$

$$A^{\text{pole}}(\Lambda_c^+\to\Sigma^+\pi^0)=\left[\frac{\sqrt{2}h^+_{\Sigma^+\Lambda_c^+}}{f_\pi}\right]\left[1-(M_{\Lambda_c^+}-M_{\Sigma^+})\left[\frac{(1+x)}{2(M_{\Lambda_c^+}-M^*_{\Sigma^+})}+\frac{(1-x)}{2(M^*_{\Sigma_c^+}-M_{\Sigma^+})}\right]\right], \tag{34}$$

$$A^{\text{pole}}(\Lambda_c^+\to\Xi^0K^+)=\left[\frac{h^+_{\Sigma^+\Lambda_c^+}}{f_K}\right](M_{\Lambda_c^+}-M_{\Xi^0})\times\left[\frac{-x}{(M_{\Lambda_c^+}-M^*_{\Sigma^+})}+\frac{5+x}{6(M^*_{\Xi_c^{0A}}-M_{\Xi^0})}-\frac{1-x}{2(M^*_{\Xi_c^{0S}}-M_{\Xi^0})}-\frac{1-x}{3(M^*_{\Xi_c^{01}}-M_{\Xi^0})}\right]. \tag{35}$$

The above formulas show explicitly the commutator term and the on-shell correction to it. The quantity x has been calculated in [8,35] and was found to be $x\approx\frac{1}{15}$, a very small value. In [35] various other test wave functions were used to establish that x was smaller than $\approx\frac{1}{10}$. We will therefore set $x=0$ in our calculation. Thus we are able to determine the S-wave pole contribution completely from the commutator term and relevant masses; no new parameter is necessary. Since the mass difference ratio in the s channel is $(M_{B_c}-M_B)/(M_{B_c}-M^*_{B'})\approx 2$ and in the u channel $(M_{B_c}-M_B)/(M^*_{B'_c}-M_B)\approx 0.7$ as mentioned previously, the s-channel contribution to the on-shell correction is more important than that of the u channel. Similar expressions for Ξ_c^{0A} and Ξ_c^{+A} decays are given in the Appendix.

V. CALCULATION AND DISCUSSION

With the baryon form factors and matrix elements determined, we are now able to calculate the amplitudes for charmed-baryon decays. We will restrict ourselves to Cabibbo-favored $B_c\to B+P(0^-)$ decays of Λ_c^+, Ξ_c^{0A}, and Ξ_c^{+A}.

First, we calculate the quark decay amplitude in the factorization approximation. In the case of $\bar{K}^0$ emission, we use the dipole q^2 dependence given in (10). For π^+ emission the effect of q^2-dependence is completely negligible. The results are shown in columns 2 and 7 of Tables I and II. We also show the results corresponding to "old factorization" [$\xi=\frac{1}{3}$ in (8)] in square brackets in this column. For the pole term the intermediate ground $\frac{1}{2}^+$ and excited $\frac{1}{2}^-$ baryon states that contribute to the decays listed in Tables I and II are as follows (according to the order in the tables): $\{\Sigma^+,\Sigma_c^0\}$; $\{\Sigma^+\}$; $\{\Sigma^+,\Sigma_c^0\}$; $\{\Sigma^+,\Sigma_c^+\}$; $\{\Sigma^+,\Xi_c^{0A},\Xi_c^{0S},\Xi_c^{01}\}$; $\{\Xi^0\}$; $\{\Xi^0,\Sigma_c^+,\Lambda_c^+,\Lambda_c^{+1}\}$; $\{\Xi^0,\Sigma_c^0\}$; $\{\Xi^0,\Sigma_c^0\}$; $\{\Xi^0,\Xi_c^{0A},\Xi_c^{0S},\Xi_c^{01}\}$; $\{\Lambda_c^+,\Lambda_c^{+1},\Sigma_c^+\}$; $\{\Xi_c^{0A},\Xi_c^{0S},\Xi_c^{01}\}$. Note that SU(4) 4-plet baryons can only be excited $\frac{1}{2}^-$ states and contribute only to the S-wave pole term. For P-wave baryon poles we use the experimental masses and we set $M_{\Xi_c^{0S}}=M_{\Xi_c^{0A}}=2.47$ GeV. The masses of the two s-channel $\frac{1}{2}^-$ noncharmed baryons $\Xi^0(\frac{1}{2}^-)$ and $\Sigma^+(\frac{1}{2}^-)$ are given by $M^*_{\Sigma^+}=M_\Sigma+0.6$ GeV and $M^*_{\Xi^0}=M_\Xi+0.6$ GeV. The mass of the excited $\frac{1}{2}^-$ charmed baryon for the u channel is given by the corresponding $\frac{1}{2}^+$ baryon mass plus a common ΔM: $M^*_{B_c}=M_{B_c}+\Delta M$. We also set, for example, $M^*_{\Lambda_c^+}=M^*_{\Lambda_c^{+1}}$. $\Delta M\simeq 0.3$ GeV is taken from [37], but fortunately the uncertainty related to this ΔM does not affect the S-wave pole contribution very much since the u-channel contribution is not as important as that of the s channel. Tables I and II show the results of our calculation. For S waves we separately list the two parts of the pole term, the soft-limit result, i.e., the commutator, in column 3 and the on-shell correction in column 4. This is to show how strong the effect of the on-shell correction is and how it affects the total $\frac{1}{2}^-$ pole contribution. The $\frac{1}{2}^-$ pole term as a whole is listed in column 5. Finally, the sum of factorization and pole contributions gives the total amplitude. We have also presented the total amplitude obtained through the use of "old factorization" [$\xi=\frac{1}{3}$ in (8)].

TABLE I. Λ_c^+ decay amplitudes (in units of 10^{-7}). The results in the square brackets correspond to the use of "old factorization".

Decay	A^{fact}	$(A^{\text{com}}$	$A^{\text{corr}})$	A^{pole}	A^{tot}	B^{fact}	B^{pole}	B^{tot}
$\Lambda_c^+ \to \Lambda\pi^+$	-6.61	(0	3.63)	3.63	-2.97	20.90	2.20	23.10
	[-5.67]				[-2.50]	[17.94]		[20.13]
$\Lambda_c^+ \to p\bar{K}^0$	-5.48	(-4.94	12.43)	7.50	2.02	13.57	5.09	18.66
	[-1.10]				[5.34]	[2.71]		[7.80]
$\Lambda_c^+ \to \Sigma^0\pi^+$	0	(-8.57	11.76)	3.19	3.19	0	6.70	6.70
$\Lambda_c^+ \to \Sigma^+\pi^0$	0	(8.57	-11.80)	-3.21	-3.21	0	-6.67	-6.67
$\Lambda_c^+ \to \Xi^0 K^+$	0	(0	0.00)	0.00	0.00	0	9.26	9.26

As can be seen in Tables I and II, different treatments of factorization can give very different results, as pointed out earlier. Due to the smaller value of $h^+_{\Sigma^+\Lambda_c^+}$, both P-wave and S-wave pole contributions are reduced, compared to the case where $h^+_{\Sigma^+\Lambda_c^+}$ is obtained from the SU(4) relation (21) and hyperon matrix element (25). However, the pole term is still considerably large in many decays and often even larger than the factorization contribution. The on-shell correction is generally very important. It always has an opposite sign to the commutator term. This is also true in hyperon decays [8], but due to the small mass difference there, the on-shell correction could only cause a small reduction in the commutator contribution. In charmed-baryon decays, however, this correction is often larger than the soft limit and thus changes the sign of the $\frac{1}{2}^-$ pole term and often even the sign of the total S-wave amplitude. In some decays such as $\Xi_c^{+A} \to \Sigma^+\bar{K}^0$ and $\Xi_c^{+A} \to \Xi^0\pi^+$ there is only the u-channel contribution and thus the on-shell correction is relatively small. In cases where there is no commuator contribution at all, the on-shell correction is not necessarily zero, for instance in $\Lambda_c^+ \to \Lambda\pi^+$, where the on-shell contributions from s and u channels do not cancel.

The quantities that can be directly compared with experimental data are the decay rate Γ and the asymmetry parameter α. Our calculated results are shown in Tables III and IV. Results using "old factorization" are also shown. Experimental data are meager. We first look at $\Lambda_c^+ \to P\bar{K}^0$. The calculated Γ is in agreement with experiment. Theoretical models often gave values 3–5 times larger than the measured rate [13–18]. In [17,18] part of the reason is that the P-wave pole term is overestimated due to the $h^+_{\Sigma^+\Lambda_c^+}$ obtained from (21). Besides, the S-wave commutator added constructively to A^{fact} (no on-shell correction was included). Thus, a larger S-wave amplitude was also obtained. The asymmetry parameter α in this decay is sensitive to the on-shell correction. In our calculation, α has changed sign due to this correction and has become smaller in magnitude. Therefore an experimental measurement of α would be very helpful in understanding the effect of the on-shell correction in this decay.

In $\Lambda_c^+ \to \Lambda\pi^+$, there is no commutator contribution but the on-shell correction exists and partly cancels A^{fact}. Thus, α is also smaller than the one given in [17,18] but has the same sign. Our α agrees with experiment. Our $\Gamma(\Lambda_c^+ \to \Lambda\pi^+)$, although smaller than those estimated in many other calculations [17,18], still seems too large compared to the measured one [5].

In $\Lambda_c^+ \to \Sigma^0\pi^+$ there is only the contribution from the pole term. The calculated decay rate is consistent with a very recent measurement [7]. Our α is quite large and would have an opposite sign if the on-shell correction were ignored. Experimental measurement of α can directly help to understand the importance of the on-shell correction in this channel. The same is true for $\Lambda_c^+ \to \Sigma^+\pi^0$ due to isospin symmetry.

In $\Lambda_c^+ \to \Xi^0 K^+$ there is only a small P-wave pole con-

TABLE II. $\Xi_c^{(0,+)A}$ decay amplitudes (in units of 10^{-7}). The results in the square brackets correspond to the use of "old factorization".

Decay	A^{fact}	$(A^{\text{com}}$	$A^{\text{corr}})$	A^{pole}	A^{tot}	B^{fact}	B^{pole}	B^{tot}
$\Xi_c^{0A} \to \Xi^0\pi^+$	8.33	(6.06	-12.22)	-6.15	2.17	-29.35	-6.96	-36.31
	[7.15]				[0.99]	[-25.19]		[-32.15]
$\Xi_c^{0A} \to \Sigma^+ K^-$	0	(0	0)	0.00	0.00	0	-9.41	-9.41
$\Xi_c^{0A} \to \Sigma^0\bar{K}^0$	-2.65	(3.49	-2.96)	0.53	-2.12	8.28	-2.49	5.79
	[-0.5]				[0.00]	[1.66]		[-0.83]
$\Xi_c^{0A} \to \Lambda\bar{K}^0$	-1.35	(-6.05	11.31)	5.26	3.91	4.19	6.40	10.59
	[0.27]				[5.00]	[0.84]		[7.24]
$\Xi_c^{0A} \to \Xi^0\pi^0$	0	(-8.57	12.09)	3.51	3.51	0	17.32	17.32
$\Xi_c^{+A} \to \Sigma^+\bar{K}^0$	3.74	(-4.94	4.19)	-0.75	2.99	-11.80	12.90	1.10
	[0.75]				[0.00]	[-2.36]		[10.54]
$\Xi_c^{+A} \to \Xi^0\pi^+$	-8.23	(6.06	-4.81)	1.25	-6.98	29.30	-17.54	11.76
	[-7.07]				[-5.82]	[25.15]		[7.61]

TABLE III. Λ_c^+ decay rates (in units of 10^{11} s^{-1}) and asymmetry parameters. The results in the square brackets correspond to the use of "old factorization".

Decay	Γ	Γ (expt)	α	α (expt)
$\Lambda_c^+\to\Lambda\pi^+$	0.81	0.39 ± 0.17[a]	-0.67	$-1.0^{+0.4}_{-0.0}$[b], -0.96 ± 0.42[c]
	[0.59]		[-0.59]	
$\Lambda_c^+\to p\bar{K}^0$	0.60	0.84 ± 0.32[d], 1.17 ± 0.42[a]	0.51	
	[0.50]		[0.78]	
$\Lambda_c^+\to\Sigma^0\pi^+$	0.17	0.35 ± 0.20[e]	0.92	
$\Lambda_c^+\to\Sigma^+\pi^0$	0.17		0.91	
$\Lambda_c^+\to\Xi^0K^+$	0.05		0	

[a]Source: Ref. [5].
[b]Source: Ref. [6].
[c]Source: Ref. [7].
[d]Source: Ref. [1].
[e]Calculated using $B(\Lambda_c^+\to\Sigma^0\pi^+)/B(\Lambda_c^+\to\Lambda\pi^+)=(2.0\pm0.7\pm0.4)/(2.2\pm0.3\pm0.4)$ from [7] and $\Gamma(\Lambda_c^+\to\Lambda\pi^+)$ from [5], and ignoring the systematic errors in the above ratio.

tribution and thus both Γ and α are very small. The situation in Ξ_c^{0A} and Ξ_c^{0S} is similar. Strong interplay between pole and factorization contributions also shows up here. In $\Xi_c^{0A}\to\Sigma^0\bar{K}^0$, $\Xi_c^{0A}\to\Lambda\bar{K}^0$, $\Xi_c^{0A}\to\Xi^0\pi^0$, and $\Xi_c^{+A}\to\Sigma^+\bar{K}^0$ the on-shell correction causes a sign change in the S-wave amplitude and thus also in α. The P-wave pole contribution can be quite large in some channels, making it hard to tell what the dominant contribution is.

The simplest decay to describe theoretically is $\Xi_c^{0A}\to\Sigma^+K^-$ which, like $\Lambda_c^+\to\Xi^0K^+$, receives only a P-wave pole-term contribution.

From the above results we can see that the charmed-baryon decays are generally quite complicated. Both factorization and pole contributions can be comparable to each other or one dominates over the other, depending on which channel and which wave is studied. Experimental data, which exist only in a few decays, do not as yet discriminate between different theoretical schemes. Despite this, the following points can be made from our study:

(1) Among the decays shown in Tables I and II, $\Lambda_c^+\to\Xi^0K^+$ and $\Xi_c^{0A}\to\Sigma^+K^-$ are the ones where there is only a P-wave pole contribution. The asymmetry parameter α is zero or very small in these two decays. Experimental measurements would be a direct test of the pole term.

(2) The on-shell correction is very important. The asymmetry parameter α is quite sensitive to this correction, especially in $\Xi_c^{0A}\to\Xi^0\pi^0$ and $\Lambda_c^+\to\Sigma^0\pi^+$ (as well as $\Lambda_c^+\to\Sigma^+\pi^0$) decays where there is no factorization contribution. The deviation from the soft limit due to the on-shell correction directly affects the magnitude and the sign of α and thus makes these channels very good candidates to test the on-shell correction.

(3) Our numerical results also show that the factorization contribution is generally not the dominant one, especially in S-wave amplitudes. If "old factorization" is used, the pole contribution dominates in decays with $\bar{K}^0$ emission. Thus more reliable baryon form factors and Hamiltonian matrix elements are needed especially for decays where both the pole terms and factorization contribute.

In the preceding we have presented a calculation of $B_c\to B+P(0^-)$ type of charmed-baryon decays. There may be uncertainties associated with the various parameters used in our model, but there is an additional uncertainty due to the effect of final-state interactions (FSI). FSI play an important role in understanding many decays of D mesons [9,11,12]. FSI have noticeable influence in hyperon decays [8] and may be more important here. The discrepancy in our calculated Γ in $\Lambda_c^+\to\Lambda\pi^+$ could well be due to the role of FSI. On the experimental side, the few measured quantities in Λ_c^+ are not enough to provide us with a test of the role played by FSI. We therefore look forward to having more experimental data available, especially on those theoretically simple decays as pointed out above. Additionally, the two $B_c\to B(\frac{3}{2}^+)+P(0^-)$ charmed-baryon decays, $\Lambda_c^+\to\Delta^{++}K^-$ and $\Xi_c^0\to\Omega^-K^+$, are very good candidates to test the pole model. There is no factorization contribution here and the D-wave amplitude in this type of decay is strongly suppressed; only the P-wave amplitude is

TABLE IV. $\Xi_c^{(0,+)A}$ decay rates (in units of 10^{11} s^{-1}) and asymmetry parameters. The results in the square brackets correspond to the use of "old factorization".

Decay	Γ	α
$\Xi_c^{0A}\to\Xi^-\pi^+$	1.55	-0.38
	[1.18]	[-0.20]
$\Xi_c^{0A}\to\Sigma^+K^-$	0.11	0
$\Xi_c^{0A}\to\Sigma^0\bar{K}^0$	0.09	-0.99
	[0.00]	[-0.02]
$\Xi_c^{0A}\to\Lambda\bar{K}^0$	0.33	1.00
	[0.36]	[0.81]
$\Xi_c^{0A}\to\Xi^0\pi^0$	0.50	0.92
$\Xi_c^{+A}\to\Sigma^+\bar{K}^0$	0.10	0.24
	[0.13]	[0.00]
$\Xi_c^{+A}\to\Xi^0\pi^+$	0.76	-0.81
	[0.49]	[-0.69]

relevant. In fact, the former channel has already been measured [1] and the latter recently observed [38]. Furthermore, Ω_c^0, which although not yet observed, is predicted to exist and decay weakly. One decay mode of this particle, $\Omega_c^0 \to \Omega^- \pi^+$, is actually *the only possible flavor combination in all Cabibbo-favored singly-charmed baryon decays which receives a contribution only from factorization.* This is also true for $\Omega_c^0 \to \Omega^- \rho^+$ since the same flavor combination is involved. The ratio of decay rates will be free of many uncertainties, thus providing a test of factorization. We hope to be able to report on these good-candidate decays in the future.

VI. CONCLUSION

We have made a calculation of the Cabibbo-favored $B_c \to B + P(0^-)$ decays of Λ_c^+, Ξ_c^{0A}, and Ξ_c^{+A}. We have pointed out the difference between the use of "old factorization" and "new factorization" in charmed-baryon decays. We favor the latter. We have calculated the relevant baryon matrix elements of the weak Hamiltonian which directly determine the pole term. The on-shell correction has been introduced and shown to be very important. Our result is in good agreement with $\Gamma(\Lambda_c^+ \to P\bar{K}^0)$, $\Gamma(\Lambda_c^+ \to \Sigma^0 \pi^+)$, and $\alpha(\Lambda_c^+ \to \Lambda \pi^+)$ but the calculated $\Gamma(\Lambda_c^+ \to \Lambda \pi^+)$ is still larger than the data, indicating a possible role played by FSI. We have also isolated several good candidates for testing models experimentally, particularly the decays $\Omega_c^0 \to \Omega^0 \pi^+$ and $\Omega_c^0 \to \Omega^- \rho^+$. Future experimental measurement on these decays will be very helpful.

Note added. After finishing this paper, we received a paper by H. Y. Cheng and B. Tseng (IP-ASTP-1791, December, 1991) who have also studied Λ_c^+ nonleptonic decays. Though our method differs in several details, Cheng and Tseng do emphasize the use of what we have referred to as "new factorization" and do include the $\frac{1}{2}^-$ pole term in their calculation.

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APPENDIX

Here we give the formulas for S-wave $\frac{1}{2}^-$ pole amplitudes in $\Xi^{(0,+)A}$ decays:

$$A^{\text{pole}}(\Xi_c^{0A} \to \Xi^- \pi^+) = \left[\frac{h^+_{\Sigma^+\Lambda_c^+}}{f_\pi}\right] \left[1 - \frac{(M_{\Xi_c^{0A}} - M_{\Xi^-})}{(M_{\Xi_c^{0A}} - M^*_{\Xi^0})}\right], \tag{A1}$$

$$A^{\text{pole}}(\Xi_c^{0A} \to \Sigma^+ K^-) = \left[\frac{h^+_{\Sigma^+\Lambda_c^+}}{f_K}\right] (M_{\Xi_c^{0A}} - M_{\Sigma^+})$$
$$\times \left[\frac{x}{(M_{\Xi_c^{0A}} - M^*_{\Xi^0})} + \frac{1-x}{2(M^*_{\Sigma_c^+} - M_{\Sigma^+})} - \frac{5+x}{6(M^*_{\Lambda_c^+} - M_{\Sigma^+})} + \frac{1-x}{3(M^*_{\Lambda_c^{+1}} - M_{\Sigma^+})}\right], \tag{A2}$$

$$A^{\text{pole}}(\Xi_c^{0A} \to \Sigma^0 \bar{K}^0) = \left[\frac{h^+_{\Sigma^+\Lambda_c^+}}{\sqrt{2} f_K}\right] \left[1 - (M_{\Xi_c^{0A}} - M_{\Sigma^0}) \left[\frac{x}{(M_{\Xi_c^{0A}} - M^*_{\Xi^0})} + \frac{1-x}{(M^*_{\Sigma_c^0} - M_{\Sigma^0})}\right]\right], \tag{A3}$$

$$A^{\text{pole}}(\Xi_c^{0A} \to \Lambda \bar{K}^0) = \left[-\sqrt{\tfrac{3}{2}} \frac{h^+_{\Sigma^+\Lambda_c^+}}{f_K}\right] \left[1 - (M_{\Xi_c^{0A}} - M_\Lambda \left[\frac{2+x}{3(M_{\Xi_c^{0A}} - M^*_{\Xi^0})} + \frac{1-x}{3(M^*_{\Sigma_c^0} - M_\Lambda)}\right]\right], \tag{A4}$$

$$A^{\text{pole}}(\Xi_c^{0A} \to \Xi^0 \pi^0) = \left[\frac{-\sqrt{2} h^+_{\Sigma^+\Lambda_c^+}}{f_\pi}\right] \left[1 - (M_{\Xi_c^{+A}} - M_{\Xi^0}) \left[\frac{1}{2(M_{\Xi_c^{0A}} - M^*_{\Xi^0})} + \frac{5+x}{12(M^*_{\Xi_c^{0A}} - M_{\Xi^0})}\right.\right.$$
$$\left.\left. + \frac{1-x}{4(M^*_{\Xi_c^{0S}} - M_{\Xi^0})} - \frac{1-x}{6(M^*_{\Xi_c^{01}} - M_{\Xi^0})}\right]\right], \tag{A5}$$

$$A^{\text{pole}}(\Xi_c^{+A} \to \Sigma^+ \bar{K}^0) = \left[\frac{-h^+_{\Sigma^+\Lambda_c^+}}{f_K}\right]$$
$$\times \left[1 - (M_{\Xi_c^{0A}} - M_{\Sigma^+}) \left[\frac{1-x}{2(M^*_{\Sigma_c^+} - M_{\Sigma^+})} + \frac{5+x}{6(M^*_{\Lambda_c^+} - M_{\Sigma^+})} - \frac{1-x}{3(M^*_{\Lambda_c^{+1}} - M_{\Sigma^+})}\right]\right], \tag{A6}$$

$$A^{\text{pole}}(\Xi_c^{+A} \to \Xi^0 \pi^+) = \left\{ \frac{h^+_{\Sigma^+\Lambda_c^+}}{f_\pi} \right\} \left[1 - (M_{\Xi_c^{+A}} - M_{\Xi^0}) \left\{ \frac{5+x}{6(M^*_{\Xi_c^{0A}} - M_{\Xi^0})} + \frac{1-x}{2(M^*_{\Xi_c^{0S}} - M_{\Xi^0})} - \frac{1-x}{3(M^*_{\Xi_c^{01}} - M_{\Xi^0})} \right\} \right] . \tag{A7}$$

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